Ultra High Energy Neutrinos from Hidden-Sector Topological Defects

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Abstract

We study Topological Defects (TD) in hidden (mirror) matter as possible sources of ultra-high energy neutrinos. The hidden/mirror and ordinary matter are assumed to interact very weakly through gravity or superheavy particles. An inflationary scenario is outlined in which superheavy defects are formed in hidden/mirror matter (and not in ordinary matter), and at the same time the density of mirror matter produced at the end of inflation is much smaller than that of ordinary matter. Superheavy particles produced by hidden-sector TD and the products of their decays are all sterile in our world. Only mirror neutrinos oscillate into ordinary neutrinos. We show that oscillations with maximal mixing of neutrinos from both worlds are possible and that values of Δm^2, needed for for solution of solar-neutrino and atmospheric-neutrino problems, allow the oscillation of ultra-high energy neutrinos on a timescale of the age of the Universe. A model of mass-degenerate visible and mirror neutrinos with maximal mixing is constructed. Constraints on UHE

neutrino fluxes are obtained. The estimated fluxes can be 3 orders of magnitude higher than those from ordinary matter. Detection of these fluxes is briefly discussed.

I. INTRODUCTION

A hidden sector of mirror particles was first suggested by Lee and Yang [1] in 1956 to save the conservation of parity in the whole enlarged particle space. This concept was further discussed and developed in Ref. [2]. Later the idea of two weakly interacting sectors, visible and hidden, found interesting phenomenological applications and development [3,4]. It has been boosted in 1980s by superstring theories with $E_8 \times E_8'$ symmetry. The particle content and symmetry of interactions in each of the E_8 groups are identical, and thus the mirror world has naturally emerged.

The most recent reincarnation of hidden-sector models is in the context of D-branes [5,6]. In this approach, light particles are associated with the endpoints of open strings which are attached to D-branes. Ordinary and hidden-sector particles live on different branes which are embedded in a higher-dimensional compactified space.

How do the ordinary and hidden sectors communicate with each other?

Most naturally they interact gravitationaly. This possibility is employed in [7]. More generally, and this is also discussed in [7] and [8], ordinary and mirror matter can also interact through the exchange of superheavy gauge particles. In the D-brane context, in some models the interaction between different branes occurs only through the exchange of closed strings (gravitons), while other models (in which the two branes are embedded in a brane of higher dimensionality) allow for a gauge boson mediated interaction.

In the case of gravitational coupling of the two worlds, one should expect dimension 5 gravitational interaction scaled by the Planck mass m_{pl}; in the case of superheavy gauge bosons this scale might be $\Lambda < m_{pl}$:

$$\mathcal{L} \sim \frac{1}{\Lambda}(\psi H)(\psi' H'), \tag{1}$$

where ψ and H are respectively the lepton and Higgs SU(2) doublets, with mirror fields denoted by primes. Eq.(1) provides mixing of ordinary and mirror neutrinos and neutrino masses [9], [10]. Because of the smallness of the neutrino masses, this is the most visible physical effect caused by the gravitational interaction ($\Lambda = m_{pl}$) of particles from the two worlds.

In principle, there could be other ways of communication. For example, (see [12]) one can add to the Lagrangian a Higgs potential term $\lambda\phi\phi\phi'\phi'$ and a gauge boson kinetic mixing term $hF_{\mu\nu}F'_{\mu\nu}$, with λ and h being new coupling constants. These terms have potential problems.

The discrete P-symmetry that interchanges the two worlds can be spontaneously broken. In this case, the coupling constants, the Higgs potentials and expectation values, and even the symmetry breaking patterns will generally differ from one world to the other. The breaking of P-symmetry can be implemented by giving a non-zero vev to a spin-0 field which transforms as a singlet under the gauge groups in both sectors ans as a pseudoscalar under the P-transformation: $P\phi = -\phi$. Models of this type have been studied in Refs. [7,11]; we shall refer to them as *asymmetric* hidden sector models.[1]

A model with an unbroken P-symmetry has been developed in Ref. [12]. The P-transformation in this model turns the left-handed ordinary fermions into right-handed mirror fermions. The masses and couplings of ordinary and mirror particles are identical, and hence the term "mirror matter" is more justified in this case. The EW Higgs fields in both sectors are also parity partners and have equal vev's. Mixing of neutrinos from different

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[1] One can, of course, consider models without P-symmetry, in which the ordinary and mirror sectors are not symmetric to begin with. For our purposes, such models are essentially the same as models with a spontaneously broken P-symmetry.

worlds (taken as an ad hoc terms in the Lagrangian) provides a channel of communication between the worlds: ordinary neutrinos can oscillate into mirror neutrinos, which play the role of sterile neutrinos. We shall refer to models of this type (with model [12] as the most elaborated example) as *symmetric* hidden sector model.

Although the term “mirror sector” suggests that it is related to the ordinary sector by a reflection transformation, most of the discussion in this paper is applicable to a more general class of hidden-sector models. In what follows, we shall use the terms “mirror sector” and “hidden sector” interchangeably.

Mirror neutrinos are the most natural candidates for sterile neutrinos often considered now for explaining the oscillations of solar and atmospheric neutrinos. In this respect mirror neutrinos have been studied in Refs. [7,12].

Mirror matter has cosmological consequences which result in model restrictions.

In symmetric hidden sector models, the number of massless and light particles is doubled in comparison with ordinary matter, and this case is excluded by cosmological nucleosynthesis, if the temperatures of mirror and ordinary matter are the same. (More generally, in any mirror model the effective number of light degrees of freedom is larger than for the ordinary matter, and this number is restricted by nucleosynthesis.)

One way to suppress the light degrees of freedom is to diminish the temperature of the mirror matter in the Universe [12,11] (see further discussion in Section II). This reduces the number density of mirror photons in a sraightforward way, while the situation with sterile neutrinos is more delicate [13]. Even if the initial density of mirror neutrinos is negligible, they reappear again and may be brought to equilibrium at nucleosynthesis epoch due to oscillations between ordinary and mirror neutrinos [14]. Nucleosynthesis constraint bounds the allowed neutrino properties in the parameter space (Δm^2 , $\sin^2 2\theta$), where θ is the mixing angle [15]. It is clear that the larger the mixing angle is, the smaller are the allowed values of Δm^2. Electron neutrinos impose the strongest limit on Δm^2, because $\nu' \leftrightarrow \nu_e$ oscillation

influences nucleosynthesis not only through the rate of the cosmological expansion, but also due to the rate of $n \leftrightarrow p$ conversion.

A cruicial assumption involved in deriving the bounds described above is that the relic lepton asymmetry is absent. In the presense of a large lepton asymmetry, $L \gtrsim 5 \cdot 10^{-5}$, the potential for active neutrinos results in a small mixing angle between active and sterile (mirror) neutrinos, and thus in a weak oscillation between these components at temperatures relevant for nucleosynthesis [13], [16], [17]. Therefore, in this case $\nu' \leftrightarrow \nu$ oscillations and induced by them the additional number of degrees of freedom are suppressed. [2].

A cosmological origin for the temperature difference between mirror and ordinary matter in the universe has been already considered in Refs. [4], [11]. The asymmetry is generated as a result of different rates of inflaton decay to ordinary and mirror matter due to P-symmetry breaking.

In Ref. [11] mirror neutrinos and baryons were considered as dark matter particles. The problem of structure formation with mirror matter and other astrophysical implications were studied in Refs. [11], [18], [19].

In this paper we shall study mirror matter in the Universe as a source of ultra-high energy neutrinos. As concrete sources we shall consider mirror Topological Defects. They produce high energy mirror neutrinos in the usual way: through production and decay of superheavy mirror X-particles. Then high energy mirror neutrinos oscillate into ordinary neutrinos, while the other products of decay of mirror X-particles remain in the mirror world, being invisible in the ordinary matter. These sources give an ideal example of "hidden

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[2] In the scenario studied in this paper only muon neutrinos from ordinary and mirror worlds are maximally mixed, while the other two types of neutrinos might have very small mixing. Therefore, even in case $L = 0$, this scenario gives only one extra neutrino and thus marginally survives the nucleosynthesis restrictions

neutrino sources" [20]. High energy neutrino radiation from ordinary sources is usually accompanied by other radiations, most notably by high energy gamma rays. Even in cases when high energy photons are absorbed in the source, their energy is partly transformed into low energy photon radiation: X-rays or thermal radiation. The fluxes of these radiations impose an upper bound on the high energy neutrino flux. For sources transparent for HE gamma radiation, in particular for Topological Defects (TD), the upper limit on diffuse neutrino flux is imposed by the cascade e-m radiation [21], [20]. In all cases (e.g. decay of X-particles, pp and $p\gamma$ interaction) neutrinos and photons are produced at the decays of pions. Colliding with microwave photons, high energy photons and electrons produce e-m cascade with most of the energy being in the observed 100 $MeV - 10\ GeV$ energy range. The energy density of this cascade radiation should not exceed, according to EGRET observations, $\omega_{cas} \sim 1 - 2 \cdot 10^{-6}\ eV/cm^3$. Introducing the neutrino energy density for neutrinos with individual energies higher than E, $\omega_\nu(> E)$, it is easy to derive the following chain of inequalities (from left to right):

$$\omega_{cas} > \omega_\nu(> E) = \frac{4\pi}{c} \int_E^\infty E I_\nu(E) dE > \frac{4\pi}{c} E \int_E^\infty I_\nu(E) dE = \frac{4\pi}{c} E I_\nu(> E), \tag{2}$$

An upper bound on the integral diffuse neutrino flux immediately follows From Eq.(2):

$$I_\nu(> E) < \frac{c}{4\pi} \frac{\omega_{cas}}{E}. \tag{3}$$

The upper bound given by Eq.(3) is the most restrictive one for diffuse neutrino fluxes produced by "ordinary" TD, AGN, gamma-ray bursts, etc. The neutrino flux from hidden-sector TD is free from this bound, though in Section V we shall present another (weaker) upper limit on high-energy neutrino flux from hidden-sector TD.

Another problem where UHE neutrinos from mirror TD can be helpful is the production of Ultra High Energy Cosmic Rays (UHECR) by resonant neutrinos.

The signature of extragalactic UHECR, the Greisen-Zatsepin-Kuzmin (GZK) cutoff [22] is not found in observations [23]. Three ideas have been suggested to explain the absence of

the cutoff: (i) signal carriers are not absorbed on the microwave radiation, (ii) the sources form a compact group around our Galaxy with a linear size smaller than the GZK absorption lenght, and (iii) the sources are distributed uniformly in the Universe, but the target particles on which UHECR are produced by the signal carriers form a compact object.

The case (iii) is realized [24] with the help of UHE neutrinos (signal carriers) interacting with relic neutrinos which have an enhanced density in some Large Scale Structure near us. One of the problems with this idea is that it is not clear how one can arrange a large flux of UHE neutrinos. The origin of this flux is left unspecified in all publications (*e.g.* [25], [26]), [27] that we are aware of. The very high resonant neutrino energy, $E_0 \sim 4 \cdot 10^{12}\ GeV$ for $m_\nu \sim 1\ eV$, implies a top-down scenario, but the fluxes of neutrinos in such scenarios are limited by the cascade constraints and are too small to produce the observed flux of UHECR. Mirror TDs, which evade the cascade constraints, can in principle produce the desirable UHE neutrino flux. This problem is addressed in Section VI.

The ouline of this paper is as follows.

In Section II we review inflationary scenarios for the mirror Universe, elaborating in particular a two-inflaton scenario for symmetric HS models. Formation of hidden-sector TD is discussed in Section III. Fluxes of High Energy Neutrinos from TD are calculated in Section IV for the case of necklaces as an example. An upper bound on the neutrino flux from hidden-sector TD is obtained in Section V. UHECR from resonant neutrinos are studied in Section VI. The main results of this paper are summarized in Section VII. In Appendix our model for mass-degenerate neutrinos with maximal mixing is described.

Throughout the paper we shall use following abbreviations: HE and UHE for High Energy and Ultra High Energy, respectivey, UHECR - for Ultra High Energy Cosmic Rays, TD - for Topological Defects, HS - for Hidden Sector (including mirror sector), LSS - for a Large Scale Structure in the universe, LG - for the Local Group of galaxies, LS -for the Local Supercluster of galaxies.

II. INFLATIONARY SCENARIOS FOR HIDDEN-SECTOR UNIVERSE

As was discussed in the Introduction, there are several dangers to be watched for in models with mirror matter. The main one is a possible conflict with the standard nucleosynthesis. If the mirror sector contains a massless photon and three light neutrinos, and the temperatures of the two worlds are the same, then the density of mirror matter at the time of nucleosynthesis is unacceptably high (it amounts to five extra neutrino species).

Two conditions are necessary to overcome this problem (see Introduction): The mirror matter must have a lower temperature $T' \lesssim 0.5T$ and lepton asymmetry is needed to suppress excessive production of mirror neutrinos through oscillation of ordinary neutrinos. Here, we shall discuss some inflationary scenarios that can naturally lead to a temperature difference between the two worlds.

The condition $T' \lesssim 0.5T$ can be easily satisfied in asymmetric hidden sector models. For example, we could have a single inflaton field ϕ which transforms as a scalar (or pseudoscalar [7]) under the P-transformation. Since the symmetry between the two sectors is broken, the field ϕ will generally have different couplings to particles in different sectors. It will then decay into ordinary and mirror particles at different rates, and the reheating temperature in the mirror matter can be lower [4,11].

This scenario would not work in symmetric hidden sector models: the inflaton would then have identical couplings to both sectors, and the two reheating temperatures would be the same. This problem can be addressed in the following two-inflaton scenario (see also [4]).

Let us consider two inflaton fields, ϕ and ϕ', with ϕ belonging to the visible sector and ϕ' to the mirror sector. During inflation, both inflatons roll down towards the minima of their respective potentials. Inflation ends when both of them have reached their minima. An important point is that the evolution of ϕ and ϕ' need not be synchronized. The inflaton dynamics is influenced by quantum fluctuations which cause inflation to end at different

times in different regions of space [28]. In our case there are two inflatons, their fluctuations are uncorrelated, and one expects them to reach their minima at different times, even in the same spatial regions. In regions where ϕ' reaches minimum first, any mirror particles produced due to its oscillations are diluted by the remaining inflation driven by the field ϕ. By the time when the energy of ϕ thermalizes, the density of mirror matter will then be very small, so that $T \gg T'$. Note that the (co-moving) coherence length of the inflaton fields should be much greater than the present horizon, so we expect very large (super-horizon) regions of the universe to be dominated by ordinary matter, and similarly large regions dominated by mirror matter, with relatively tiny boundary regions where both kinds of matter are present in comparable amounts. It is very unlikely for us to find ourselves in one of such rare regions.

A quantitative analysis shows that this de-synchronization picture may or may not apply, depending on the form of the inflaton potential $U(\phi)$. We shall see, however, that it does apply for the simplest choice of the potential,

$$U(\phi) = m_\phi^2 \phi^2/2. \tag{4}$$

In "chaotic" inflation scenario, the inflatons roll from very large values of ϕ towards $\phi = 0$. The initial values of ϕ and ϕ' are large and uncorrelated. At very large ϕ, $\phi > \phi_q$, the dynamics of ϕ are dominated by quantum fluctuations. The boundary ϕ_q of this "quantum diffusion" regime is determined by condition $dU/d\phi \sim H^3$. At $\phi_* \lesssim \phi \lesssim \phi_q$, where $\phi_* \sim m_{pl}$, ϕ (and ϕ') evolve in a slow-roll regime described by the equations

$$3H\dot{\phi} = -\frac{dU(\phi)}{d\phi}, \tag{5}$$

$$3H\dot{\phi}' = -\frac{dU(\phi')}{d\phi'}, \tag{6}$$

$$H^2 = \frac{8\pi G}{3}[U(\phi) + U(\phi')], \tag{7}$$

Will the "incidental" initial ratio ϕ'/ϕ be conserved during the slow-roll evolution? For the potential (4) the answer is "yes". Indeed, the integration of Eqs. (5) - (7) results in

$$\phi'/\phi = const. \tag{8}$$

Suppose that $\phi' \ll \phi$ at the time when ϕ begins its slow roll ($\phi \sim \phi_q \sim m_{pl}^{3/2} m_\phi^{-1/2}$), so that we can neglect $U(\phi')$ in (7). The characteristic value of ϕ' is then determined by quantum fluctuations about $\phi' = 0$ and can be found from [29] $U(\phi') \sim H_q^4$, where $H_q \sim (m_\phi m_{pl})^{1/2}$ is the expansion rate at $\phi \sim \phi_q$. This gives $\phi' \sim m_{pl}$ and $\phi'/\phi \sim (m_\phi/m_{pl})^{1/2}$. This small ratio of ϕ'/ϕ is preserved all the way to the end of inflation. In this type of models, the universe is divided into super-horizon regions dominated by ordinary matter and equally large regions dominated by mirror matter.

For a different choice of the potential, e.g.

$$U(\phi) = \lambda_\phi m_{pl}^4 (\phi/m_{pl})^n, \tag{9}$$

with $n > 2$, the fields ϕ and ϕ' do get synchronized at late stages of the evolution. (For $n = 4$, this was shown in Ref. [30].) In this case, integration of the slow-roll equations (5) -(7) gives

$$\phi^{-n+2} - \phi'^{-n+2} = const. \tag{10}$$

At the onset of the slow roll of ϕ, when $\phi \sim \phi_q \sim \lambda^{-1/4} m_{pl}$, the typical value of ϕ' is $\phi' \sim \lambda_\phi^{(2-n)/n(n+2)} m_{pl}$. With $\lambda_\phi \ll 1$, this is much smaller than ϕ_q but still much larger than ϕ_*. Now, consider the solution (10) with these initial values of ϕ and ϕ'. By the end of inflation, both fields get much smaller than their initial values, so that the constant on the right-hand side of (10) becomes unimportant, and we have $\phi' \approx \phi$. Thus, ϕ and ϕ' get synchronized by the end of inflation, even if they were not initially. We conclude that models with a power-law potential (9) and $n > 2$ give equal densities of mirror and ordinary matter.

Coming back to our basic ordinary-matter dominated scenario, we can give another example of a model with segregated mirror and ordinary matter is a two-inflaton model where inflation occurs at a metastable minimum of the inflaton potential. The highest rate of inflation is achieved in the false vacuum state where both ϕ and ϕ' are at the minima of their respective potentials. This state decays through nucleation of two types of bubbles. In bubbles of the first type, the field ϕ' tunnels through the barrier and starts rolling down its potential, while ϕ remains in the false vacuum. As ϕ' rolls to the bottom of the potential and decays into mirror particles, inflation continues in the interior of the bubble. Mirror particles are quickly diluted away, and the bubble interior is filled with inflating false vacuum of the "ordinary" inflaton ϕ. This vacuum will in turn decay through nucleation of bubbles of the second type (with ϕ tunneling and ϕ' remaining unchanged). If the potential is sufficiently flat, the roll down of ϕ is accompanied by additional inflation, and after the ϕ-field decay, the interiors of these secondary bubbles will become dominated by ordinary matter. Quite similarly, nucleation of ϕ'-bubbles inside ϕ-bubbles results in regions dominated by mirror matter.

III. HIDDEN-SECTOR TOPOLOGICAL DEFECTS

Apart from the general considerations which apply to any model with a hidden sector (HS), we have to address some additional issues specific to HS defects. First we have to arrange for these defects to form. And second, we have to avoid the formation of similar defects in our sector, since otherwise we would get unacceptably large fluxes of ordinary UHE particles, and the cascade bound would be violated.

Once again, these conditions can be easily satisfied in asymmetric HS models, where the symmetry breaking scales and even the symmetry breaking patterns may be different in the two worlds. Perhaps the simplest possibility is the model of Ref.([11]) with one inflaton and asymmetric reheating. HS defects can be formed in a usual symmetry-breaking phase

transition after inflation. We only have to arrange for the corresponding phase transition in the ordinary matter to occur at a lower energy scale or not to occur at all. [3]

In symmetric HS models, the two sectors have identical physics and identical defect solutions. It does not follow, however, that ordinary and HS defects should be present in the universe in equal numbers. The density of defects is determined by the cosmological evolution, which can be different for the two sectors. For example, in the two-inflaton model discussed above, mirror matter is completely inflated away, and if defects were formed in phase transitions after inflation, we would expect to have ordinary defects but no HS defects. For our purposes, however, we need the opposite situation: HS defects and no ordinary defects.

This can be arranged if defects are formed in a curvature-driven phase transition during inflation.

As an illustration we shall consider a toy model of a spontaneously broken $SU(2)$ symmetry. We introduce a Higgs triplet $\chi = (\chi_1, \chi_2, \chi_3)$ with a potential

$$V(\chi) = \frac{1}{4}\lambda(|\chi|^2 - \eta^2)^2. \tag{11}$$

When χ acquires an expectation value, $SU(2)$ is broken to $U(1)$ and monopoles are formed. Suppose now that χ is coupled to the inflaton ϕ,

$$V_\phi(\chi) = -\frac{1}{2}g\phi^2|\chi|^2, \tag{12}$$

and has a non-minimal coupling to spacetime curvature,

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[3] In models discussed in Refs. [7,11], the P-symmetry breaking is assumed to be at the electroweak scale, $\eta_P \sim \eta_{EW} \sim 10^2 GeV$, and the two worlds have nearly identical physics above η_{EW}. But this needs not be the case. The physics of mirror and ordinary defects will be different in models where η_P is greater than the symmetry breaking scale of the defects.

$$V_R(\chi) = \frac{1}{2}\xi R|\chi|^2. \tag{13}$$

The mirror field χ' has identical couplings to the mirror inflaton ϕ' and to the curvature.

We shall assume "chaotic" inflation with a quadratic inflaton potential (4). After the mirror inflaton ϕ' reached its minimum, the effective mass of the mirror field χ' is

$$m_{\chi'}^2 = -\lambda\eta^2 + \xi R, \tag{14}$$

and the curvature is given by

$$R \approx 12H^2 \approx 16\pi(m_\phi/m_{pl})^2\phi^2. \tag{15}$$

We see that above the critical curvature

$$R_{\chi'} = \lambda\eta^2/\xi. \tag{16}$$

the expectation value of χ' vanishes and the $SU(2)'$ symmetry is restored. Thus, even in the absence of mirror matter, as the curvature gradually decreases during inflation, we can have symmetry-breaking phase transitions accompanied by the formation of HS defects. If these curvature-driven phase transitions occur sufficiently close to the end of inflation, the defects will not be completely inflated away and can serve as sources of UHE neutrinos.

The effective mass of the field χ in the ordinary sector is

$$m_\chi^2 = -\lambda\eta^2 - g\phi^2 + \xi R = -\lambda\eta^2 - g[1 - 16\pi(\xi/g)(m_\phi/m_{pl})^2]\phi^2. \tag{17}$$

We see that $m_\chi^2 < 0$ for all ϕ, provided that

$$16\pi(\xi/g)(m_\phi/m_{pl})^2 < 1. \tag{18}$$

In this case no ordinary defects are formed during the whole slow roll period of ϕ (and any defects formed prior to that are completely diluted away)

One final condition that has to be checked is that the temperature at reheating is not so high that the symmetry gets restored, since otherwise "ordinary" defects will be formed again in a subsequent phase transition. All the above conditions can be satisfied without fine-tuning.

IV. FLUXES OF HIGH-ENERGY NEUTRINOS FROM HS DEFECTS

We shall consider necklaces [32] as a specific example of TD. Necklaces are hybrid TDs formed by monopoles (M) and antimonopoles ($\bar{M}$), each being attached to two strings. The monopole mass m and the mass per unit length of string μ are determined by the corresponding symmetry breaking scales , η_s and η_m,

$$m \sim 4\pi\eta_m/e, \qquad \mu \sim 2\pi\eta_s^2 \tag{19}$$

where e is the gauge coupling. The evolution of necklaces depends on the parameter

$$r = m/\mu d \tag{20}$$

which gives the ratio of the monopole mass to the average mass of string between two monopoles (d is the average string length between monopoles). It cannot exceed $r_{max} \sim \eta_m/\eta_s$. As it is argued in ref. [32], necklaces might evolve towards a scaling solution with a constant $r \gg 1$, possibly approaching $r \sim r_{max}$. Monopoles and antimonopoles trapped in the necklaces inevitably annihilate in the end, producing superheavy Higgs and gauge bosons (X particles) of mass $m_X \sim e\eta_m$. The rate of X-particle production per unit volume and time is

$$\dot{n}_X \sim r^2\mu/t^3 m_X \tag{21}$$

From the relations above it is easy to see that

$$\zeta \equiv \frac{r^2\mu}{m_X^2} = \frac{2\pi}{e^2}\left(\frac{r}{r_{max}}\right)^2 \lesssim 10. \tag{22}$$

High energy neutrinos are produced in the chain of X-decays via pions. For simplicity we assume that pions (of all charges) are produced with a power-law spectrum

$$D_\pi(x, m_X) = 4(2-p)2^{-p}x^{-p} \tag{23}$$

where $x = E_\pi/m_X$ is a fraction of energy taken away by a pion, and for p we shall use a value between 1.3 and 1.5 , which bound the realistic QCD spectrum of pions. The spectrum (23) is normalized so that $\int_0^{1/2} x D_\pi(x, m_X) dx = 1$.

The number of neutrinos with energy E_ν from the decay of one X-particle is given by

$$N(E_\nu) = \frac{4}{m_X} \int_{2E_\nu/m_X}^{1/2} \frac{dx}{x} D_\pi(x, m_X). \tag{24}$$

The diffuse flux of mirror neutrinos ν_i' (where $i = e$ and μ , antineutrinos are not included) is

$$I_{\nu_i}(E) = \frac{c}{\pi} \frac{\dot{n}_X(t_0)}{m_X H_0} \int_0^\infty dz \int_{2E_\nu(1+z)/m_X}^{1/2} \frac{dx}{x} D_\pi(x, m_X) \tag{25}$$

Finally, we obtain the diffuse flux of ordinary neutrinos ν_i taking into account $\nu_i' \to \nu_i$ oscillation with averaged probability $P_{osc} \sim 1/2$:

$$I_{\nu_i}(E) = \frac{c}{4\pi} \frac{\dot{n}_X(t_0)}{m_X H_0} \left(\frac{E_\nu}{m_X/2} \right)^{-p} \frac{(2-p)2^{2-p}}{p(p-1)} P_{osc} \tag{26}$$

This expression can be written in more compact form:

$$I_{\nu_i} = k_p \zeta \frac{c}{4\pi} \frac{1}{t_0^2} \left(\frac{E_\nu}{m_X} \right)^{-p}, \tag{27}$$

with ζ given by Eq.(22) and

$$k_p = \frac{6(2-p)2^{-2p}}{p(p-1)} P_{osc}. \tag{28}$$

For $P_{osc} = 1/2$ and $p = 1.5$ $k_p = 1/4$ and $k_p = 0.89$ for $p = 1.3$.

The neutrino flux in Eq.(27) can be very large. For example, with $r^2\mu \sim 0.1 m_X^2$, $m_X \sim 10^{16}\ GeV$, $p = 1.5$ and $P_{osc} = 1/2$, one obtains at $E \sim 10^{11}\ GeV$, $E^3 I_\nu \sim 1 \cdot 10^{28}\ eV^2 m^{-2} s^{-1} sr^{-1}$, i.e. a flux three orders of magnitude larger than predicted from ordinary sources under most optimistic assumptions. Since we have used a very rough power-law approximation for the spectrum, it may be better to illustrate the enhancement of the flux

by comparing Eq. (26) to the maximum allowed flux from ordinary sources with the same power-law spectrum.

Suppose that some unidentified ordinary-matter sources produce superheavy X-particles, which decay producing high energy pions. Then the energy density of neutrinos ω_ν and the cascade density ω_{cas} are equal. Using this fact we immediately obtain an upper limit on the neutrino flux as

$$I_{\nu_i}(E) \le (2-p)2^{-p}\frac{\omega_{cas}}{m_X^2}\frac{c}{4\pi}\left(\frac{E}{m_X}\right)^{-p} \tag{29}$$

The ratio of the two fluxes [Eq.(26) and Eq.(29)] is given basically by the value $m_X^2/(\omega_{cas}t_0^2)$ and it is $\sim 2\cdot 10^4$ for $m_X \sim 1\cdot 10^{16}$ GeV and for observationally allowed cascade density $\omega_{cas} \sim 2\cdot 10^{-6}$ eV/cm^3.

V. E-M CASCADE RESTRICTIONS

All particles from the QCD cascades produced by decays of mirror X-particles are sterile in our world. Only mirror neutrinos can oscillate into ordinary ones. An upper bound on the neutrino flux is given by the resonant interaction of UHE neutrinos with relic cosmological neutrinos, $\nu + \bar{\nu} \to Z^0 \to pions$. Electrons and photons from the decay of pions initiate e-m cascades on the microwave radiation. Reactions $\nu + \bar{\nu} \to Z^0 \to l + \bar{l}$, where $l = e, \mu, \tau$, also contribute to the cascade. The calculated cascade energy density ω_{cas} must be smaller than the energy density ω_γ^{obs} observed (*e.g.* by EGRET) in the extragalactic diffuse radiation.

Let us first derive a convenient formula for the rate of resonant events.

The resonant neutrino energy E_0 and the resonant cross-section $\sigma(E)$ for $\nu+\bar{\nu} \to Z^0 \to f$ (f is an arbitrary final state) are given by

$$E_0 = \frac{m_Z^2}{2m_\nu} = 4.16\cdot 10^{12}\left(\frac{1\ eV}{m_\nu}\right)\ GeV \tag{30}$$

$$\sigma_{\nu,f}(E_c) = \frac{12\pi}{m_Z^2}\frac{\Gamma_\nu\Gamma_f}{(E_c - m_Z)^2 + \Gamma_t^2/4} \tag{31}$$

Here, m_Z is the mass of Z^0-boson, E_c is the center-of-mass energy, Γ_ν, Γ_f and Γ_t are the widths of Z^0 decay to neutrinos, to an arbitrary final state f and the total width, respectively. In Eq.(31) we took into account that only one chiral component of neutrino takes part in the interaction. [4]

The rate of Z^0 production per unit volume due to collisions of high energy flavor neutrino ν_i with target antineutrino $\bar{\nu}_i$ is given by

$$\dot{n}_Z = 4\pi n_{\bar{\nu}_i} \xi \int I_{\nu_i}(E)\sigma(E)dE = 4\pi\sigma_t n_{\bar{\nu}_i} \xi I_{\nu_i}(E_0)E_0, \tag{32}$$

where

$$\sigma_t = 48\pi f_\nu G_F = 1.29 \cdot 10^{-32}\ cm^2, \tag{33}$$

G_F is the Fermi constant, and here and below $f_s = \Gamma_s / G_F m_Z^3$, with Γ_s being the width of the channel. Numerically, $f_\nu = 0.019$, $f_h = 0.197$ and $f_{tot} = 0.283$. The case of HE $\bar{\nu}_i$ can be trivially added. Summation over i takes place in the case of mass-degenerate neutrinos. For the target neutrino density we shall use one helicity density with zero chemical potential, $n_{\nu_i} = 56\ cm^{-3}$, corresponding to the temperature $T = 2.73K$ of the microwave radiation, and $\xi = \cos^2\theta$, $\sin^2\theta$ or 1, as explained in the footnote.

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[4]Production of Z^0 occurs through the interaction of flavor neutrinos, *e.g.* $\nu_{eL} + \bar{\nu}_{eL}$ in the case of Dirac neutrinos, or $\nu_{eL} + \nu_{eL}^c$ in the case of Majorana neutrinos. In practice one considers interaction of HE flavor neutrino, *e.g.*, $\bar{\nu}_e$, with a physical mass-eigenstate target neutrino ν_1 of mass m_1. The probability to find this neutrino as ν_e is equal to $\cos^2\theta$ (or $\sin^2\theta$), where θ is the mixing angle. Therefore, the cross-section in Eq.(31) must include $\xi = \cos^2\theta$ (or $\sin^2\theta$). In the case of mass-degenerate neutrinos, $m_1 \approx m_2 = m_\nu$, an incident HE $\bar{\nu}_e$ interacts with ν_e component of ν_2 almost at the same resonant energy and $\xi = 1$. There is no difference in the rates for Dirac and Majorana neutrinos: this becomes obvious if in counting the number of neutrino species one includes both ν_L and ν_L^c in the case of Majorana neutrinos.

Note that Eqs.(32),(33) are exact formulae because integration in Eq.(32) takes place over a narrow resonant peak.

The rate of Z^0 production with subsequent decay of Z^0 to an arbitrary channel f is given by $\dot{n}_Z(Z \to f) = n_Z f_f/f_{tot}$. Taking into account only dominant hadron channels, with $f_h/f_{tot} = 0.696$, and using the fact that pions transfer to electro-magnetic cascade half of their energy, one can find the energy density of electromagnetic cascade as

$$\omega_{cas} = \frac{1}{2}\frac{f_h}{f_{tot}} E_0 \dot{n}_Z \xi t_0 = 2\pi\xi\sigma_t n_{\nu_i} t_0 I_{\nu_i}(E_0) E_0^2 \frac{f_h}{f_{tot}}. \quad (34)$$

Using $\omega_{cas} \leq \omega_\gamma^{obs}$ we obtain an upper limit

$$I_{\nu_i}(E_0) \leq \frac{2\omega_\gamma^{obs} m_\nu^2}{\pi\sigma_t n_{\nu_i} t_0 m_Z^4 \xi} \frac{f_{tot}}{f_h} \quad (35)$$

One can see from Eq.(35) that unless the production spectrum of neutrinos has a cutoff at some energy lower than E_0, the upper limit is provided by the lightest relic neutrino ν_i. This is not surprising: in QCD spectra most of the energy is concentrated in particles of the highest energies. The lower m_ν is, the higher is E_0 and the more energy it tranferred to the e-m cascade.

The scale of neutrino masses suggested by oscillation solutions to the atmospheric neutrino and solar neutrino problems is $m \sim \sqrt{\Delta m^2}$. This gives the following masses: $m_\nu \approx 5 \cdot 10^{-2}\ eV$ from atmospheric neutrino oscillations, $m_\nu \approx 2 \cdot 10^{-3}\ eV$ from SMA MSW, $m_\nu \approx 4 \cdot 10^{-3}\ eV$ from LMA MSW and $m_\nu \approx 9 \cdot 10^{-6}\ eV$ from VO solution. The latter mass requires a value of m_X which is too large, so we disregard this case.

It is easy to verify that we did not exceed the e-m cascade limit in the calculations of section IV. Indeed, Eqs.(26) and (34) with $r^2\mu/m_X^2 = 0.1$, $m_X = 1 \cdot 10^{16}\ GeV$, $n_{\nu_i} = 56\ cm^{-3}$, with E_0 given by Eq.(30) and with both channels $\nu_i + \bar{\nu}_i$ and $\bar{\nu}_i + \nu_i$ taken into account, result in $\omega_{cas} = 1.2 \cdot 10^{-7}/\sqrt{m_1}\ eV/cm^3$, where $m_1 = m_\nu/1\ eV$. The energy density ω_{cas} is well below the allowed limit for $m_1 = 1$ and is marginally below it for $m_1 = 3 \cdot 10^{-3}$.

VI. UHECR FROM RESONANT NEUTRINOS

Here we shall estimate the Ultra High Energy Cosmic Ray (UHECR) fluxes produced by neutrinos from hidden-sector TD. If the target neutrino density is enhanced in nearby large-scale structures (LSS), such as the halo (h) of our Galaxy, the Local Group (LG) and the Local Supercluster (LS), the large flux of the observed UHECR could be generated there. Photon fluxes dominate in the production spectra. The ratio of photon fluxes from a large-scale structure, $I_\gamma^{LSS}(E)$, and from extragalactic space, $I_\gamma^{extr}(E)$, can be expressed in terms of the overdensity of target neutrinos in the structure, δ_{LSS}, and the length of gamma-ray absorption in extragalactic space, $R_\gamma(E)$, as

$$I_\gamma^{LSS}(E)/I_\gamma^{extr}(E) = \delta_{LSS} R_{LSS}/R_\gamma(E), \tag{36}$$

where R_{LSS} is the linear size of the large-scale structure.

The overdensity factors for the galactic halo (h), Local Group (LG), and Local Supercluster (LS) are discussed in the accompanying paper [31], and for non-degenerate neutrinos they are estimated as

$$\delta_\nu^h < 37 m_1^3, \qquad \delta_\nu^{LG} < 13 m_1^3. \qquad \delta_\nu^{LS} \sim 1. \tag{37}$$

Here, $m_1 = m_\nu/1\ eV$ is the neutrino mass in units of 1 eV and the overdensity is defined as the ratio of the neutrino density with flavor i (antineutrinos are not included) in the structure to the average density of the same neutrinos in the Universe, $n_{\nu_i} = 56\ cm^{-3}$. From Eqs.(36),(37) one can see that while LS does not give an enhancement of UHECR flux, both the galactic halo ($R_h \sim 100\ kpc$), and the Local Group ($R_{LG} \sim 1\ Mpc$) give an enhancement of order $(0.3-1)m_1^3$. Note that this excess flux arrives without absorption. Estimates for both structures are given as upper limits, with the limit for the galactic halo being more reliable. In the estimates below, the index LSS refers to one of these two structures.

Once again, let us take HS necklaces as an example of neutrino sources, with the diffuse neutrino flux $I_{\nu_i}(E)$ given by Eq.(27). Using the formalism developed in Section V, one can

write down the flux of Z^0-bosons with resonant energy E_0:

$$I_{Z^0} = 2\xi\sigma_t\delta_\nu n_{\nu_i} R_{LSS} I_{\nu_i}(E_0)E_0.$$

Assuming a power-law spectrum of hadrons in $Z^0 \to hadrons$ decay and using Eq.(27) for $I_{\nu_i}(E_0)$, one can easily calculate the UHE photon flux from the halo or LG:

$$I_\gamma^{LLS}(E) = k_\gamma \xi\zeta\sigma_t\delta_\nu n_{\nu_i} R_{LSS} P_{osc}\frac{c}{4\pi t_0^2}\left(\frac{E}{m_X}\right)^{-p}, \tag{38}$$

where R_{LSS} is $\sim 1\ Mpc$ is the case of LG, and $\sim 100\ kpc$ in the case of galactic halo, $\zeta = r^2\mu/m_X^2$, $\xi = \cos^2\theta$, $\sin^2\theta$ or 1 and k_γ is given by

$$k_\gamma = \frac{4(2-p)^2 2^{-2p}}{p(p-1)}\frac{\Gamma_{had}}{\Gamma_{tot}} \tag{39}$$

where $\Gamma_{had}/\Gamma_{tot}$ is the ratio of Z^0 decay widths, equal to 0.7. For $p = 1.5$ $k_\gamma = 0.116$ and for $p = 1.3$ $k_\gamma = 0.58$.

With $\delta_\nu = \delta_\nu^{max} \sim m_\nu^3$, the UHE gamma-ray flux given by Eq.(38) is proportional to m_ν^3 and m_X^p. For a fixed overdensity, the flux does not depend on the neutrino mass and depends only on the mass of X-particle as m_X^p.

As a numerical example let us consider the case of a gamma-ray flux from LG with two neutrino flavors and with a degenerate mass $m_\nu = 2\ eV$ ($\xi = 1$), taking the mass of X-particle $m_X = 1\cdot 10^{15}\ GeV$ and $\zeta = \zeta_{max} = 10$. For $p = 1.5$ and $p = 1.3$, the values of $E^3 I_\gamma(E)$ at $E = 1\cdot 10^{20}\ eV$ are equal to $2.3\cdot 10^{24}\ eV^2/m^2ssr$ and $1.8\cdot 10^{24}\ eV^2/m^2ssr$, respectively, i.e. close to the observed values.

It is interesting to derive an upper limit for the UHE gamma-ray flux inside a LSS, using e-m cascade production in the space outside it. For LSS with a linear size R_{LSS} and neutrino overdensity δ_ν, one obtains

$$I_\gamma^{max}(E) = \frac{2}{3}(2-p)\frac{\delta_\nu R_{LSS}}{ct_0}\frac{c}{4\pi}\frac{\omega_{cas}}{E_0^2}\left(\frac{E}{E_0}\right)^{-p}. \tag{40}$$

From Eq.(40) one can see that the upper limit does not depend on m_X and is proportional to m_ν^{5-p} for $\delta_\nu = \delta_\nu^{max}$. For the parameters of LG ($R_{LG} = 1\ Mpc$ and $\delta_\nu^{LG} = 13m_1^3$)

and $\omega_{cas} = 1 \cdot 10^{-6}\ eV/cm^3$ one obtains, at $E = 1 \cdot 10^{20}\ eV$, fluxes equal to $E^3 I_\gamma(E) = 5.0 \cdot 10^{23} m_1^{3.5}\ eV^2/m^2 ssr$ and $2.4 \cdot 10^{23} m_1^{3.7} eV^2/m^2 ssr$ for $p = 1.5$ and $p = 1.3$, respectively. At $m_\nu > 2\ eV$ both upper limits are consistent with observations.

Turning the argument around, one can obtain a lower limit on the neutrino mass from the condition $I_\gamma^{max}(E) > I_{obs}(E)$ at $E = 1 \cdot 10^{20}\ eV$: $m_\nu \gtrsim 2\ eV$.

More accurate calculations with realistic QCD spectra from Z^0 decay are given in the accompanying paper [31]

VII. DISCUSSION AND CONCLUSIONS

Mirror matter is a natural option in models with $G \times G'$ symmetry, in particular in superstring models $E_8 \times E_8'$. The coupling constants, the Higgs vev's, and the symmetry breaking patterns in the two sectors may or may not be the same, depending on whether or not the discrete P-symmetry interchanging the sectors is spontaneously broken.

We assume that the two sectors communicate due to a non-renormalizable interaction (1), where the case $\Lambda \sim m_{pl}$ corresponds to gravitational interaction. These interactions result in neutrino masses and neutrino oscillations, including the oscillations between ordinary and mirror (sterile) neutrinos.

Cosmological restrictions rule out a wide class of hidden-sector models; they are particularly severe for the symmetric models in which the P-symmetry is unbroken. In such models, the number of light particles is doubled, and this is excluded by the cosmological nucleosynthesis. The nucleosynthesis constraints can be avoided by suppressing the temperature of the mirror matter, accompanied by an introduction of a lepton asymmetry, which suppresses $\nu \to \nu'$ oscillation.

Inflationary scenarios resulting in different temperatures in the ordinary and mirror sectors can easily be constructed for asymmetric HS models. In the case of symmetric models, we discussed a two-inflaton scenario, first outlined in [4]. The inflatons ϕ and ϕ' belong to the

ordinary and hidden sector, respectively. In regions of space where ϕ' reaches the minimum of its potential earlier than ϕ, the products of ϕ' decay are diluted by the expansion driven by the ordinary inflaton ϕ. When ϕ also rolls to the bottom of its potential, we get a superhorizon region dominated by ordinary matter. In stochastic inflation, superbubbles dominated by mirror matter are equaly often produced. We have shown that this scenario can work only with a suitable choice of the inflaton potential: for some choices the slow rolls of the two inflatons get synchronized, resulting in equal densities of mirror and ordinary matter. We also suggested an alternative version of the two-inflaton model where the potential has a metastable minimum. Then the inflating false vacuum decays through nucleation of ϕ- and ϕ'-bubbles, and the segregation of ordinary and mirror matter is achieved in a natural way.

Despite the suppression of mirror matter, hidden-sector topological defects can dominate over the ordinary ones. Once again, this can be easily arranged in asymmetric HS models. In symmetric models, the two sectors have the same types of defects with identical properties, but the cosmological densities of the defects need not be the same. We illustrated this possibility by a two-inflaton model with a curvature-driven phase transition (see section III). In this model, HS topological defects are produced in a phase transition during inflation, when the mirror inflaton ϕ' is already at the minimum of its potential. The phase transition is triggered when the spacetime curvature (which is driven by the ordinary inflaton potential) decreases to some critical value. If this happens sufficiently close to the end of inflation, the resulting defects are not inflated away. The corresponding phase transition in the ordinary matter occurs much earlier, and ordinary topological defects are completely diluted by inflation. Thus, in the two-inflaton scenario we can have a desirable situation when the universe is dominated by ordinary matter and hidden-sector topological defects.

HS topological defects produce high-energy neutrinos in the chain of decays of superheavy particles – constituent fields of the defects. All decay products are invisible in the ordinary matter and only mirror neutrinos oscillate into the ordinary world. The flux of neutrinos

from ordinary topological defects is limited by cascade photons which are produced in the same decays of pions as neutrinos. This restriction is absent in the case of HS defects. However, the cascade limit for mirror neutrinos reappears, though in a weaker form. After $\nu' \to \nu$ oscillation, ordinary neutrinos produce hadrons, e^+e^- $\mu^+\mu^-$, and $\tau^+\tau^-$ in the resonant scattering off the background (dark matter) neutrinos: $\bar{\nu} + \nu_{DM} \to Z^0 \to hadrons$, or l^+l^-. These particles (or products of their decay) initiate electromagnetic cascades on the microwave photons. The smaller is the mass of DM neutrinos, the stronger is the cascade upper limit [see Eq.(35)]. This is because the resonant energy is inversely proportional to the neutrino mass: $E_0 = m_Z^2/2m_\nu$. In QCD spectra, most of the energy is carried by high energy particles, and thus more energy is transferred to the e-m cascade when E_0 is large.

As a specific example of mirror topological defects, sources of high energy neutrinos, we studied the necklaces – magnetic monopoles connected by strings, with each monopole being attached to two strings. We found that, for a reasonable choice of model parameters, the diffuse neutrino flux can be three orders of magnitude higher than that from ordinary necklaces, being still consistent with the cascade upper limit imposed by the resonant production of Z^0 bosons. Note, however, that the accuracy of our calculations is limited by the power-law approximation of the energy spectra.

A diffuse flux of UHE neutrinos from HS topological defects can produce the observed flux of UHE cosmic rays due to resonant interaction with the Dark Matter neutrinos in the Local Group, if the mass of the target neutrino is $m_\nu > 2\ eV$. The atmospheric-shower producing particles in this case are UHE photons. Their spectrum does not exibit a cutoff, because of the relatively small size ($R \sim 1\ Mpc$) of LG.

This model of UHE cosmic rays requires mass degenerate neutrinos with $m_\nu > 2\ eV$. Our model for neutrino masses and mixing is described in the Appendix. Mirror muon neutrino ν'_μ is maximally mixed with the ordinary ν_μ neutrino, and both have masses $m_\nu \sim 2\ eV$. Their mass difference, $\Delta m^2 \approx 2 \cdot 10^{-3}\ eV^2$, is responsible for the atmospheric neutrino

oscillations observed in Super-Kamiokande. Solar neutrino anomaly is explained by $\nu_e \to \nu_\mu$ LMA MSW solution with $\Delta m^2 \approx 4 \cdot 10^{-5}\ eV^2$ and $\sin^2 2\theta \approx 0.80$. Thus, we assume that the three neutrinos, ν_e, ν_μ and ν'_μ are maximally mixed and mass degenerate ($m_\nu \sim 2\ eV$).

The calculated neutrino flux is below the upper limits obtained from horizontal air shower observations at EAS TOP [33] and AGASA [34] at $10^6 - 10^7\ GeV$ and marginally below Fly's Eye limit [35] at $10^{11}\ GeV$. The predicted neutrino fluxes can be detected by this technique with the help of these and future bigger arrays, like e.g. "Auger" detector [36]. However, the best hope for detecting these neutrinos probably rests with the future satellite detectors such as OWL (Orbiting Wide Field Light Collector) [37] and AIRWATCH [38].

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APPENDIX A: MASS DEGENERATE NEUTRINOS IN MIRROR MODELS

We consider $G \times G'$ model with EW symmetry $SU_2(L) \times U_1$ in G and $SU'_2(R) \times U'_1$ in G'. G and G' representations communicate through operators of dimension $d = 5$ with a scale $\Lambda < m_{pl}$.

The particle content of the EW group relevant to the neutrino masses is

$$\psi_L = \begin{pmatrix} \nu_L \\ l_L \end{pmatrix}, \quad l_R, \quad \phi = \begin{pmatrix} \phi_+ \\ \phi_0 \end{pmatrix}, \quad \phi^c = \begin{pmatrix} \phi_0^* \\ -\phi_+^* \end{pmatrix} \tag{A1}$$

for $SU_2(L) \times U_1$, and

$$\psi'_R = \begin{pmatrix} \nu'_R \\ l'_R \end{pmatrix}, \quad l'_L, \quad \phi' = \begin{pmatrix} \phi'_+ \\ \phi'_0 \end{pmatrix}, \quad \phi'^c = \begin{pmatrix} \phi'^*_0 \\ -\phi'^*_+ \end{pmatrix} \tag{A2}$$

for $SU'_2(R) \times U'_1$. Here, ϕ and ϕ' are the Higgs fields and l and l' are charged leptons.

There are no light singlets ν_R and ν'_L in our model. There are 4 neutrino states: ν_L, ν^c_L, ν'_R, ν'^c_R, in terms of which the most general expression for the mass matrix is

$$\mathcal{L} \sim \left(\bar{\nu}_L \;\; \bar{\nu'}^c_R\right) \begin{pmatrix} m_L & M \\ M & m_R \end{pmatrix} \begin{pmatrix} \nu^c_L \\ \nu'_R \end{pmatrix} + h.c. \tag{A3}$$

Since $\bar{\nu'}^c_R \nu^c_L = \bar{\nu}_L \nu'_R$, there are three independent mass operators in Eq.(A3): $\bar{\nu}_L \nu'_R$, $\bar{\nu}_L \nu^c_L$, $\bar{\nu'}^c_R \nu'_R$, and they are generated in our model. Indeed, these operators are:

$$SU_2(L) \times U_1 : \quad \left(\bar{\psi}_L \phi^c\right) \left(\phi^c \psi^c_L\right) \rightarrow \bar{\nu}_L \nu^c_L \tag{A4}$$

$$SU'_2(R) \times U'_1 : \quad \left(\bar{\psi}'_R \phi'^c\right) \left(\phi'^c \psi'^c_R\right) \rightarrow \bar{\nu}'_R \nu'^c_R \tag{A5}$$

$$intergroup : \quad \left(\bar{\psi}_L \phi^c\right) \left(\phi'^{c+} \psi'_R\right) \rightarrow \bar{\nu}_L \nu'_R \tag{A6}$$

Arrows show the neutrino mass operators after the EW symmetry breaking, $< \phi_0 >= v_{EW}$ and $< \phi'_0 >= v'_{EW}$. Communication of visible and mirror sectors is accomplished by the *intergroup* term in Eq.(A6). It has a dimensional scale Λ, and thus one obtains

$$M = v_{EW} v'_{EW} / \Lambda. \tag{A7}$$

This is the basic neutrino mass scale in our model, and we want it to be $M \sim 1\ eV$. For $v'_{EW}/v_{EW} \sim 10$, we need $\Lambda \sim 10^{14}\ GeV$.

The scales of d=5 terms operating inside $SU_2(L) \times U_1$ and $SU'_2(R) \times U'_1$ groups, Λ_L and Λ_R, can be different. For our model (see below) we need the following hierachy of masses:

$$m_L < m_R \ll M. \tag{A8}$$

It can be provided by $\Lambda \ll \Lambda_R < \Lambda_L$, or in the model-dependent way. One can observe, for example, that both intragroup terms (A4) and (A5) violate the lepton number defined for

the doublets as $L_\psi = L_{\psi'} = 1$, while the intergroup operator (A6) conserves it. One can build a model with one universal Λ where the intragroup d=5 operators are forbidden and thus m_L and m_R are suppressed.

Let us assume a local $\tilde{U}_1$ symmetry for massless particles before symmetry breaking, with the following charge assignment: $q = +1$ for $\psi_L,\ \ l_R,\ \ \psi'_R,\ \ l'_L$ and $q = 0$ for $\phi,\ \ \phi'$.

The terms (A4) and (A5) do not conserve q. Let us introduce two new scalar SU_2 singlets, Φ and Φ', with $q = 2$. Now, apart from the operator (A6), we can write two other SU_2-singlet operators conserving the charge q and the electric charge. They are the following $d = 6$ operators:

$$\frac{\Phi}{\Lambda^2}\left(\bar{\psi}_L \phi^c\right)(\phi^c \psi^c_L)\,, \quad \frac{\Phi'}{\Lambda^2}\left(\bar{\psi}'_R \phi'^c\right)(\phi'^c \psi'^c_R)\,. \tag{A9}$$

After EW and $\tilde{U}_1$ symmetry breaking with vev's $<\Phi>= V$ and $<\Phi'>= V'$, respectively, one obtains

$$M = v_{EW} v'_{EW}/\Lambda,\ m_R = v'^2_{EW} V'/\Lambda^2,\ m_L = v^2_{EW} V/\Lambda^2, \tag{A10}$$

which satisfy the hierarchy (A8).

Let us now come back to the mass matrix (A3). Its diagonalization gives the masses of eigenstates and the mixing angle for visible and mirror neutrinos:

$$m_{1,2} = \left(m_R + m_L \pm \sqrt{4M^2 + (m_R - m_L)^2}\right) \approx 2M \tag{A11}$$

$$\Delta m^2 = m_2^2 - m_1^2 \approx 2(m_R + m_L)M \tag{A12}$$

$$\sin 2\theta = \frac{2M}{\sqrt{4M^2 + (m_R + m_L)^2}} \approx 1 \tag{A13}$$

The hierarchy of masses (A8) provides mass degenerate neutrinos with (almost) maximal mixing.

Till now we considered asymmetric mirror models. In the case of symmetric models [12], the masses $M,\ \ m_L$ and m_R are considered as free parameters, and thus the hierarchy condition (A8) can be arbitrarily fulfilled.

As a realistic example we can consider the case when $\nu_\mu \to \nu'_\mu$ oscillations explain the atmospheric neutrino anomaly ($\Delta m^2 \approx 2 \cdot 10^{-3}\ eV^2$) and LMA MSW explains the solar neutrino problem ($\nu_e \to \nu_\mu$ oscillation). In this case, $2M \sim 1\ eV$ and $m_R \sim 1 \cdot 10^{-3}\ eV$. All three neutrinos, ν_μ, ν_e and ν'_μ, must have mass $2M \sim 1\ eV$ and be bimaximally mixed. The oscillation length ($\nu'_\mu \to \nu_\mu$) for HE neutrinos with resonant energy E_0 is only $l_{osc} = 4\pi E_0/\Delta m^2 = 5{\cdot}10^{20}\ cm$. The values of Δm^2 in this case are $2{\cdot}10^{-3}\ eV^2$ and $\sim 4{\cdot}10^{-5}\ eV^2$ for the atmospheric and solar neutrinos, respectively; they are not affected much by radiative corrections [39]. Radiative splitting affects mostly τ neutrino, because of the large Higgs coupling with the tau. This effect is not important in our scenario where only three other neutrinos, ν_μ, ν'_μ and ν_τ, are degenerate in mass.

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